ON THE POSSIBLE USE OF INTENSE BEAMS OF THE BIG PROTON ACCELERATORS FOR EXCITATION OF A LINEAR ACCELERATOR STRUCTURE

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This work belongs to the mainstream of research pursued by the above Institute during the recent decade, and which was initiated by G. Budker. The principal idea is to accelerate particles by means of the intense beams of ultrarelativistic charged particles that are available in modern accelerators. Since these beams have a large amount of energy stored per unit length*), an efficient transfer of this energy into the electromagnetic field of a linear accelerator structure**) could significantly enhance the accelerating gradient in linacs and open up new possibilities in high-energy physics[1].

In this paper we consider a method of driving an UHF waveguide structure by using an intense ultrarelativistic beam with small emittance and small energy spread, as obtained in modern high-energy cyclic accelerators and storage rings, to generate an accelerating electric field of ultimately high intensity. We describe the arrangement of the accelerator structure, the strong-focusing guiding field lattice, and schemes of bunching. We have studied the capability of electron and proton beams existing in various high-energy laboratories, and find that in principle there is a possibility of obtaining an average accelerating gradient of $\gtrsim$ 1 MeV/cm with the use of existing machines. Finally, we discuss the potential applications of such linear accelerators.

1. ESTIMATES OF THE INDUCED FIELD STRENGTH

Let us consider the excitation of an E_{010} mode electromagnetic field in a cylindric cavity with radius a and length h. The cavity is penetrated along its axis by a beam with a current density $\vec{j}(\vec{r},t)$ that comprises N charged particles bunched in a long periodic sequence of k identical bunches such that their frequency is ω_{010}/n (n is integer). We assume that these weakly coupled cavities are in fact the cells of the linear accelerator structure (iris-loaded waveguide) into which the properly bunched accelerated beam is injected past the exciting beam. Energy conservation gives:

$$\frac{1}{8\pi}\int \vec{E}^2\, dV = \frac{1}{2}\int dt \int \eta\vec{j}(\vec{r},t)\vec{E}\, dV \; . \tag{1}$$

Here the integration is done over the cavity volume V; $\vec{E}$ is the electric field at $t \gg T = 2\pi/\omega_{010}$; η stands for the fraction of the beam fitting in the cavity aperture. We thus obtain for the amplitude E of the E_{010} mode:

$$E = \left[\frac{2}{J_1(\alpha_{01})}\right]^2 \frac{Ne}{a^2}\, \eta\, k_t\, k_f\, k_b\, k_r \; , \tag{2}$$

*) For example, the CERN Super Proton Synchrotron (SPS) 400 GeV proton beam with 3×10^{13} ppp ($\sim$ 0.2 A) corresponds to about 2 million joules of stored energy and to some 100 GW of power in a one-turn extraction pulse.

**) Obviously, the smaller is the operation wavelength of the structure (and respectively the smaller is the volume occupied by the UHF field into which the energy stored in the beam should be supplied), the higher is the acceleration field intensity: $E \propto \lambda^{-2}$.

where

- $J_m(x)$ is the first kind Bessel function of the m^{th} order, $\alpha_{m\ell}$ is the ℓ^{th} zero of J_m;
- the transit-time factor

$$k_t = \sin \frac{\pi h}{\beta\lambda} \Big/ \frac{\pi h}{\beta\lambda} ,$$

with $h \approx a$, and $\beta \equiv v/c \approx 1$, is close to unity;

- the form factor k_f is a characteristic of the beam coupling with the electric field of the mode in question:

$$k_f = \int \vec{j}\vec{E}\, dV/j_{max}EV < 1 ;$$

- the bunching factor k_b is actually the amplitude of the Fourier harmonic of j/j_{max} at the frequency ω_{010} ($k_b < 1$);
- the last factor k_r accounts for damping in the cavity:

$$k_r = \frac{1}{k} \frac{1 - \exp(-knT/\tau)}{1 - \exp(-nT/\tau)} < 1 ;$$

here τ stands for the mode damping time.

For a narrow ($\eta = k_f = 1$) completely bunched ($k_b = 1$) beam with the pulse duration $knT \ll \tau$ ($k_r = 1$), Eq. (2) yields:

$$E = E_1 \equiv \left[\frac{2}{J_1(\alpha_{01})}\right]^2 \frac{Ne}{a^2} k_t \simeq 14.9 \frac{Ne}{a^2} k_t . \qquad (3)$$

Using the relation: $\omega_{010} = c(\alpha_{01}/a)$, we rewrite E_1 in terms of λ:

$$E_1 = \left[\frac{4\pi}{\alpha_{01}\, J_1(\alpha_{01})}\right]^2 \frac{Ne}{\lambda^2} k_t \simeq 102 \frac{Ne}{\lambda^2} k_t . \qquad (4)$$

For practical estimates we express E_1 in MV/cm:

$$E_1 \left(\frac{MV}{cm}\right) = 1.47 \frac{10^{-11}N}{\lambda^2(cm)} k_t . \qquad (4')$$

In the case where the damping is essential: $knT \gtrsim \tau$ (but, of course, $nT \ll \tau$), we have

$$k_r = \frac{\tau}{knT} [1 - \exp(-knT/\tau)] ,$$

and rewrite Eq. (4):

$$E = E_1 k_r = 102 \frac{I\tau}{\lambda^2} [1 - \exp(-knT/\tau)]\, k_t . \qquad (5)$$

Here $I = Ne/knT$ is the average current. For a long beam $knT \gg \tau$, and the stationary field amplitude is completely determined by $I\tau/\lambda^2$ *). Using the relation between τ and λ,

*) If the current I is not high enough to achieve the value of the required E, it is possible to use the preliminary compression of the proton beam (in a single or multi-bunch mode) by the rotation through a proper angle in the longitudinal motion phase space.

$$\tau \approx \frac{\alpha_{01}}{2\pi s}\left(\frac{\sigma\lambda^3}{c^3}\right)^{1/2} \propto \lambda^{3/2} , \tag{6}$$

we obtain in this limiting case $E \propto \lambda^{-1/2}$. Equation (6) assumes that $h \approx a$; s is a wall smoothness factor ($s \gtrsim 1$), σ is the wall conductivity. Introducing the unit length shunt impedance,

$$R = 51\,\frac{\alpha_{01}}{2\pi s}\sqrt{\frac{\sigma}{\lambda c^3}}\,k_t \simeq 2.25\,\frac{k_t}{s\sqrt{\lambda(\mathrm{cm})}} \quad \left[\frac{\mathrm{M\Omega}}{\mathrm{cm}}\right] ,$$

we express the stationary amplitude in the form:

$$E = 2I\cdot R \simeq 4.5\,\frac{I(\mathrm{A})}{s\sqrt{\lambda(\mathrm{cm})}}\,\eta\,k_t\,k_b \quad \left[\frac{\mathrm{MV}}{\mathrm{cm}}\right] . \tag{7}$$

However, at short λ the above consideration is no longer valid since the transverse beam size is no longer negligible in comparison with the cavity diameter. The field generation efficiency is then decreased owing to the reduction of the fundamental mode electric field averaged over the beam cross-section ($k_f < 1$), and also because of a smaller beam transmission $\eta(\lambda) < 1$. In this case the outer part of the beam cross-section is of no use and should be eliminated with a proper cut-off (or it can be separated and used as an ordinary proton beam of lower intensity).

2. STRONG-FOCUSING LATTICE TO IMPROVE BEAM TRANSMISSION

It has been shown in Section 1 that in order to achieve an amplitude E that is as high as possible, the wavelength λ has to be reduced to the limit given by complete beam transmission. It is therefore advantageous to put the accelerator section into a strong-focusing magneto-optical lattice which retains small lateral beam size over the whole section length. For such an accelerator section we shall determine the wavelength λ that yields the maximum intensity E for a specified emittance $\varepsilon_x \cdot \varepsilon_z$ of the exciting beam.

To calculate the beam transmission factor $\eta(\lambda)$, we consider a model of the FODO lattice consisting of thin quadrupole lenses with optical strength P spaced by drift lengths L. We find for the advance μ of the betatron phase per period:

$$\cos\mu = 1 - \frac{1}{2}\left(LP\right)^2 , \tag{8}$$

and for the maximum of the β-function:

$$\beta_{max} = \frac{2}{P}\sqrt{\frac{2 + LP}{2 - LP}} , \tag{9}$$

and choose the operating point such as to ensure the stable motion of decelerated particles with energies between E_{min} and the initial energy $E_0 \equiv eH\rho$, and also the stability of the accelerated beam in the energy range between E_0 and E_{max} ($E_{max} \gg E_{min}$):

$$\mu = \frac{\pi}{2}\,\frac{E_{min}}{E_{max}} . \tag{10}$$

Having expressed P in terms of the short lens length L_ℓ and the magnetic field H_0 at its aperture radius a (which is assumed to be equal to that of the cavity)*), and using Eqs. (8) and (10), we obtain an equation for L_ℓ:

$$\mu \simeq LP = L\,\frac{H_0 L_\ell}{a\,H\rho} = \frac{\pi}{2}\,\frac{E_{min}}{E_{max}}\;; \tag{11}$$

β_{max} is lower if $L_\ell \approx L$. Then Eqs. (9) and (11) yield

$$\beta_{max} = 2\,\frac{aH\rho}{H_0}\sqrt{\frac{2}{\pi}\,\frac{E_{max}}{E_{min}}\,\frac{H_0}{aH\rho}}\;.$$

Thus the accepted phase area A of the accelerator section within the strong focusing lattice may be written in the form:

$$A = \frac{d^2}{4\beta_{max}} = \frac{d^2}{4}\left(\frac{\pi}{8}\,\frac{E_{min}}{E_{max}}\,\frac{H_0}{H\rho\;a}\right)^{\frac{1}{2}}\;, \tag{12}$$

where $d \equiv 2\xi a$ is the iris diameter. With the relation $a = \lambda(\alpha_{01}/2\pi)$ we have

$$A(\lambda) = \left[\frac{\pi}{8}\,\frac{E_{min}}{E_{max}}\,\frac{H_0}{H_\rho}\left(\frac{\alpha_{01}}{2\pi}\right)^3 \xi^4\,\lambda^3\right]^{1/2}\;. \tag{13}$$

If a Gaussian density distribution is assumed in the exciting beam, and if the energy spread in the beam does not contribute to its lateral size (i.e. the dispersion function $\psi \equiv 0$ along the section), then the beam transmission factor is readily obtained:

$$\eta(\lambda) = \left\{1 - \exp\left[-\frac{A(\lambda)}{2\varepsilon_x}\right]\right\}\left\{1 - \exp\left[-\frac{A(\lambda)}{2\varepsilon_z}\right]\right\}\;. \tag{14}$$

Here $A(\lambda)$ is given by Eq. (13). In the case $A \ll \varepsilon_{x,z}$ we have

$$\eta(\lambda) \simeq \frac{A^2}{4\varepsilon_x\varepsilon_z} \propto \lambda^3\;, \tag{15}$$

and substituting Eq. (15) into Eq. (4) it can be seen that in this case of a large emittance $E \propto \lambda$, and the phase-space density of the beam $N/\varepsilon_x\varepsilon_z$ plays a decisive role.

Assuming for the sake of simplicity that $\varepsilon_z \approx \varepsilon_x = \varepsilon$, we substitute Eq. (14) into Eq. (5), where k_t and k_g are taken as fixed and $\tau \to \infty$ (short beam pulses), and obtain

$$E = 102\,\frac{Ne}{\lambda^2}\left\{1 - \exp\left[-\frac{A(\lambda)}{2\varepsilon}\right]\right\}^2 = 102\,\frac{Ne}{B^2}\,\Phi\!\left(\frac{\lambda}{B}\right)\;. \tag{16}$$

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*) Note that in a case where the highest possible accelerating field E does not condition any magnetic field in the accelerating structure, the focusing lenses could be placed only between the accelerating sections, and consequently the focusing gradient would be augmented according to the shorter lenses lengths, while the aperture may be made as small as the iris size.

In the above equation we have introduced a dimensionless function $\Phi(x) = [1 - \exp(-x^{3/2})]^2/x^2$ and

$$B = \left[\frac{\pi}{32\varepsilon^2}\frac{E_{min}}{E_{max}}\frac{H_0}{H\rho}\left(\frac{\alpha_{01}}{2\pi}\right)^3 \xi^4\right]^{-1/3} . \tag{17}$$

The maximum of $\Phi(x)$, $\Phi_{max} = 0.4086$, is reached by an optimum choice of λ:

$$\lambda_{opt} = 0.835\ B \simeq 1.33\left[\frac{E_{max}}{E_{min}}\frac{E_0(\mathrm{GeV})}{H_0(\mathrm{kG})}(5\xi)^{-4}\varepsilon^2\ (\mu\mathrm{rad\cdot m})\right]^{1/3} ; \tag{18}$$

at λ_{opt},

$$E_{max} = 0.286\left(102\ \frac{Ne}{\lambda^2_{opt}}\right) . \tag{19}$$

In the opposite limiting case for a long-pulse beam, the accelerator structure approaches c.w. operation and the stationary amplitude is found by substitution of Eq. (14) into Eq. (7):

$$E\left(\frac{\mathrm{MV}}{\mathrm{cm}}\right) = 4.5\ \frac{I(A)}{s\sqrt{B}}\ F\left(\frac{\lambda}{B}\right) . \tag{20}$$

A dimensionless function $F(x) \equiv [1 - \exp(-x^{3/2})]^2/x^{1/2}$ reaches its maximum value $F_{max} = 0.6262$ under the condition

$$\lambda_{opt} = 2.044B = 3.26\left[\frac{E_{max}}{E_{min}}\frac{E_0(\mathrm{GeV})}{H_0(\mathrm{kG})}(5\xi)^{-4}\varepsilon^2\ (\mu\mathrm{rad\cdot m})\right]^{1/3} , \tag{21}$$

and the relevant maximum electric field intensity in this stationary c.w. regime is

$$E_{max} = 0.894\left[4.5\ \frac{I(A)}{s\sqrt{\lambda(\mathrm{cm})}}\right]\left[\frac{\mathrm{MV}}{\mathrm{cm}}\right] . \tag{22}$$

Note that the first numerical factor in Eqs. (19) and (22) corresponds to the optimum value of the beam transmission $\eta(\lambda_{opt})$ in the two limiting cases, respectively. Figure 1 shows the dependence of E on λ in the two cases in question [Eqs. (16) and (20)].

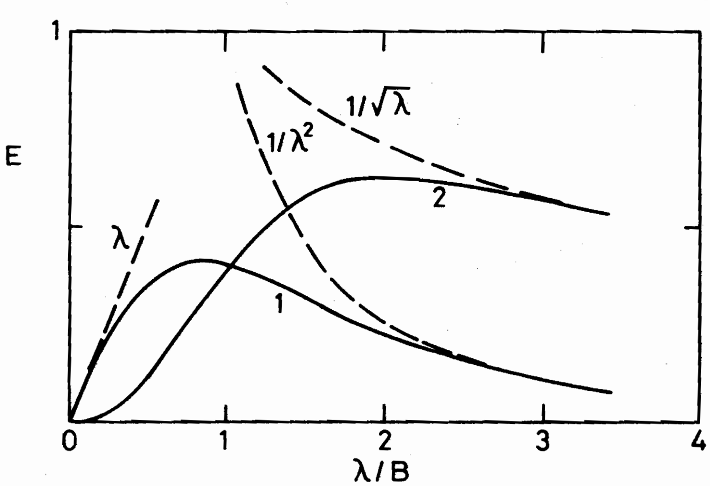

Fig. 1

The induced electric field amplitude E (arb. units) versus operation wavelength λ. Curve 1 - short bunch; curve 2 - long bunch [Eqs. (16) and (20), respectively].

To give illustrative examples for the two limiting cases Eqs. (19) and (22), we take the parameters of the 1.5 GeV electron beam extracted from the storage ring VEPP-3 (Novosibirsk) and those of the 400 GeV proton beam of the SPS (CERN). The results are summarized in Table 1.

Table 1

Summary of the numbers relevant to the problem of high accelerating gradient in a beam-driven linac

		VEPP-3	SPS
Energy	GeV	1.5	400
Beam current	A	0.3 [4)]	0.24 [2,3)]
Intensity	Particles pp	5×10^{11}	3×10^{13} [2)]
Emittance	mm•mrad	0.3	0.017 [3)]
Energy spread	MeV	1.5	40 [3)]
λ_{opt}	cm	0.6	0.9
E	MV/cm	5.7	0.75

3. CALCULATION OF THE BEAM SIZE

To justify the choice of the strong-focusing lattice parameters based on the over-simplified treatment of Section 2 [Eqs. (8) to (12)], we now study in more detail the motion of the decelerated and accelerated particles in this magneto-optical system. Let us consider the usual linearized equations of the transverse particle oscillations in paraxial approximation:

$$\frac{d}{ds}\left(p\,\frac{dz}{ds}\right) - p\,\frac{G}{H\rho}\,z = 0 \,, \tag{23}$$

$$\frac{d}{ds}\left(p\,\frac{dx}{ds}\right) + p\,\frac{G}{H\rho}\,x = 0 \,. \tag{23'}$$

Here s is the longitudinal coordinate, z and x stand for the particle lateral excursions, G is the magnetic field gradient in the lenses, and the momentum p(s) is assumed to be a specified function of s. Having written the solutions of Eqs. (23) and (23′) in the Floquet form,

$$z = \sqrt{\frac{\varepsilon_0 p_0}{p(s)}}\; w(s)\,\cos\chi(s) \,,$$

we obtain equations for the betatron phase $\chi(s)$ and for the envelope function w(s):

$$\chi' = \frac{1}{w^2} \,, \tag{24}$$

$$w'' + \left(\frac{1}{4}\frac{p'^2}{p^2} - \frac{1}{2}\frac{p''}{p} - \frac{G}{H\rho}\right) w = \frac{1}{w^3} \,. \tag{25}$$

An equation for the horizontal envelope can be obtained from Eq. (25) by the replacement: $G \to -G$.

Thus, by finding a numerical solution of Eq. (25) with initial conditions that describe the incident beam and with a specified p(s), and then computing the transverse size of the beam that is travelling through the accelerator section,

$$\left(\overline{A^2_{z,x}}\right)^{\frac{1}{2}} = \left[\frac{\varepsilon_0 p_0}{p(s)}\right]^{\frac{1}{2}} w(s) \ , \tag{26}$$

one can make an optimum choice of the strong-focusing lattice parameters such that the beam dimensions corresponding to the accelerated (decelerated) particles in all RF phases of the induced field do not exceed the aperture of the waveguide iris. [Note that for a high enough accelerating field, the gradient of the UHF magnetic field in the E_{010} mode and the iris focusing -- both of which are dependent on the RF phase -- should be accounted for in Eqs. (23) and (25) contributing in G.]

4. CALCULATION OF THE BUNCHING EFFICIENCY

For bunching of the ultrarelativistic beam with the period close to $\lambda_{opt} \sim 1$ cm, a klystron principle can be applied: after an energy modulation by an external voltage U = = $E_0 u \sin \omega t$, the beam passes through a bending magnet with the field H(s) wherein the path length differs by Δs when the particle energy differs by $\varepsilon = \Delta E/E_0$:

$$\Delta s = \int_0^s \varepsilon \psi(s') \frac{H(s')}{H\rho} ds' \equiv \varepsilon L \ . \tag{27}$$

This results in a modulation of linear charge density $\kappa(s)$. Assuming the primary beam to be Gaussian in the energy phase plane,

$$\kappa(s) = \int_{-\infty}^{\infty} \frac{1}{2\pi\sigma_\varepsilon \sigma_s} \exp\left(-\frac{\varepsilon_0^2}{2\sigma_\varepsilon^2} - \frac{s_0^2}{2\sigma_s^2}\right) d\varepsilon \ , \tag{28}$$

$$\varepsilon_0 = \varepsilon - u \sin \frac{2\pi}{\lambda} (s - \varepsilon L) \ , \tag{29}$$

$$s_0 = s - \varepsilon L \ . \tag{30}$$

Assuming $\sigma_s \gg \lambda$, we obtain a periodic $\kappa(s)$, expandable in Fourier series:

$$\kappa(s) = \sum_{n=-\infty}^{\infty} \kappa_n \exp\left(-in \frac{2\pi}{\lambda} s\right) \ .$$

The Fourier-harmonic amplitudes have the same meaning as the "bunching factor" k_b introduced in Section 1; they are determined by

$$\kappa_n = J_n\left(n \frac{u}{\sigma_\varepsilon} \frac{2\pi}{\lambda} \sigma_\varepsilon \cdot L\right) \exp\left[-\frac{1}{2}\left(n \frac{2\pi}{\lambda} \sigma_\varepsilon L\right)^2\right] \ . \tag{31}$$

The maximum value $\kappa_1 = 0.582$ is reached at $\sigma_\varepsilon = 0$ and $2\pi u L = 1.84\ \lambda$. The energy spread in the beam $\sigma_\varepsilon \neq 0$ causes a reduction in κ_n; however, the reduction is inessential in practice if $u/\sigma_\varepsilon \gtrsim 3$. For the SPS 400 GeV proton beam $\sigma_\varepsilon \lesssim 10^{-4}$, which comes to some 30 MeV in absolute units, and the modulator voltage should exceed 100 MV.

Thus the "klystron" technique readily provides the bunching factor $k_b \approx 0.5$. An enhancement of k_b can be achieved by the repeated use of the transformation (29) and (30). In the theory of multi-resonator klystrons it is generally known that even two-stage cascade bunching with an optimum choice of the parameters u and L, and of phase shift, allows us to reach $k_b \approx 0.78$ ($\sigma_\varepsilon = 0$). However, at ultrarelativistic energies such a multi-stage cascade bunching seems to be inconvenient as it requires additional drift sections, i.e. additional bending magnet arrays of large curvature radius.

Another way of enhancing the bunching efficiency is provided by an admixture of higher harmonics in the beam modulator voltage. In an idealistic case, the sawtooth-shaped modulator voltage with the normalized amplitude $u = U_{peak}/E_0$ yields

$$\kappa_n = \frac{\sin (2\pi n/\lambda)[(\lambda/2) - uL]}{(2\pi n/\lambda)[(\lambda/2) - uL]} \times \exp\left[-\frac{1}{2}\left(n\frac{2\pi}{\lambda}\sigma_\varepsilon L\right)^2\right], \tag{32}$$

and a maximum $k_b = \kappa_{1max} = 1$ at $\sigma_\varepsilon = 0$. An optimum approximation of the sawtooth by three harmonics can raise k_b to ~ 0.8. With more harmonics the approximation is known to converge slowly, and the further enhancement of k_b is not significant.

Note that the synchrotron ring itself can be used as a bending magnet for beam density modulation if the beam, after energy modulation, is made to travel around a part of a whole turn or several turns before the final beam ejection. The energy modulator structure could be placed in a synchrotron straight section or it could be placed in a special bypass, the proton beam passing through it only once before ejection. (Such a solution looks like being the cheapest one.)

5. ON NEW POSSIBILITIES IN HIGH-ENERGY PHYSICS

In the last section we will discuss some of the new possibilities in high-energy physics that are opened up by the proposed technique of a linear accelerator excited by an intense proton beam.

5.1 Acceleration of protons

Obviously, in the first stage the technique would be to accelerate some fraction of the initial proton beam that is shifted in accelerating phase. In this way it is possible to have a proton beam of doubled energy (roughly speaking) with an intensity of about one tenth of that of the primary beam. Using several such accelerating sections in series, excited by separated fractions of the initial beam, the final energy can be increased proportionally as the proton intensity decreases.

5.2 Acceleration of secondaries

A very important application of such a linear accelerator excited by a high-energy proton beam, is the acceleration of very intense beams of unstable secondary particles up to the energy of the protons, with good emittance and monochromaticity (if good phasing).

For the unstable particles not to decay during acceleration, it is necessary to have an energy gain per $c\tau_0$ of the particle (τ_0 is the rest-frame lifetime) that is more than two times larger than m_0c^2 (m_0: the rest-frame mass); this means that

$$E \gtrsim 2\,\frac{m_0 c}{\tau_0}\,. \tag{33}$$

For muons this means E > 3.2 keV/cm; for $\pi^{\pm}$, E > 0.36 MeV/cm; for $K^{\pm}$, E > 2.8 MeV/cm.

5.3 Acceleration of pions

The most difficult problem in the case of pions is to design a converter system that collects, with an efficiency of about 10%, several-GeV pions produced by hundred-GeV protons in a sufficiently small transverse emittance. There are possibilities of solving this problem by using a multi-sectioned target and parabolic and/or lithium lenses. Perhaps it will be necessary to use a special initial accelerating section.

Some words about the phasing problem. The required initial phase distribution in the pion beam can be obtained by using a properly bunched initial proton beam. It is more difficult, however, to maintain a good phasing during acceleration in cases where the difference between the velocities of the exciting and the accelerated particles is not small (for rather low-energy secondary particles, for example). In this case, after an accelerating section of the length $\frac{1}{2}\,\lambda\gamma^2_{min}$, it is necessary to install a wiggler magnet section that shifts the phase of the low-energy accelerated particles relative to the exciting particles by about $\frac{3}{2}\,\pi$. To keep the pion beam free from muons, it is necessary to accelerate it as quickly as possible, thus preventing the pions from decaying. For removing protons of the same momentum from the positive pion beam, it is better to use the difference in their velocities -- during the acceleration the above-mentioned phasing efforts give the dephasing effect in the accompanying proton motion (a kind of RF separator).

5.4 Accumulation and acceleration of muons

For producing an intense, very low emittance (and pure) muon beam, it is necessary to proceed as follows: i) collect as many multi-GeV pions as possible in the smallest phase-space volume using a hadron cascade; ii) let them decay in a very strong focusing channel (to prevent an increase of emittance); iii) cool the muons in a special storage ring or linear accelerator using ionization energy loss in a target placed in a very low β-value region, and compensating the average energy loss[5]; iv) bunch the ejected muons properly; v) finally, accelerate them in the main linear accelerator.

5.5 Acceleration of kaons?

The charged kaons acceleration will become feasible when it will be possible to obtain accelerating fields of about 3 MeV/cm or more. And here the only crucial problem is the breakdown limitation of the electric field on the resonator cavity surface, because the required peak proton current can be obtained even now, using preliminary phase compression as mentioned above. For purifying the beam, it is possible to use the above-mentioned "RF-separator" trick.

5.6 Stored, cooled, and polarized beams

Of course, it will be possible to accelerate previously stored and cooled beams in such an accelerator. This approach seems feasible for antiprotons, electrons, positrons, and ions. In the case of electrons, light ions and, maybe, positrons, the average intensity would be limited by the intensity of the main proton accelerator (to about 10% of the proton intensity). For antiprotons and heavy ions the average intensity will be limited by the production efficiency of the beams.

It is necessary to mention especially the possibility of accelerating polarized beams in such a linear accelerator without loss in the degree of polarization. The available intensity of polarized proton and electron beams will probably reach the energy conservation limit ($\sim$ 10% of proton beam intensity) in the near future, while the intensity of polarized positron and antiproton beams will be limited by their production efficiency.

5.7 Generation of neutrino beams

From the pions, muons, and kaons accelerated in such a way, it would be possible to obtain intense and well-collimated full-energy neutrino beams ($\nu_{e,\mu}$; $\bar{\nu}_{e,\mu}$). Especially good results could be obtained by using a special high-field storage-decay ring (for muons and even for pions).

5.8 Exotic colliding beams

This technique opens new possibilities for solving the problem of superhigh-energy colliding beams of unstable particles and electrons, namely: $\pi^{\pm}p$, $\pi^{\pm}\pi^{\pm}$, $\mu^{\pm}\mu^{\pm}$, $e^{\pm}e^{\pm}$ colliding beams[5,6].

For pion-pion (±) beams accelerated in the linear structure, it is necessary to build a special high-field storage ring, to inject into it as many pions as possible at a time, and to repeat the cycle with all the protons available from the main synchrotron. In this case the average pion-pion luminosity ($L_{\Sigma}^{\pi\pi}$) will be equal to

$$L_{\Sigma}^{\pi\pi} = \frac{\zeta\dot{N}_p}{\ell_t^{eff}} \frac{N_\pi}{\ell_b} \frac{p_\pi p}{(m_\pi c)^2} \frac{eH\tau_\pi}{2\pi m_\pi c} ,$$

where

ζ = the efficiency of proton-pion conversion,

$\dot{N}_p$ = number of protons per second,

N_π = the number of pions in one super-bunch*),

ℓ_t^{eff} = the effective length of the optimized conversion target,

ℓ_b = length of the pion super-bunch assumed to be equal to the β-value in the interaction region,

p_π = the pion momentum of conversion,

p = the pion momentum in colliding beams,

H = the storage ring magnetic field,

τ_π = the rest frame pion lifetime.

If we put $\dot{N}_p = 10^{13}$ pps, $N_\pi = 10^{11}$ πpp, $\zeta = 10^{-1}$, p_π = 5 GeV/c, p = 500 GeV/c, H = 100 kG, ℓ_t^{eff} = 1 cm, ℓ_b = 1 m, we have

$$L_{\Sigma}^{\pi\pi} = 3 \times 10^{27} \text{ cm}^{-2} \text{ s}^{-1} .$$

*) Actually, in a colliding beam facility, maximization of N_π may force us to choose a greater operation wavelength λ so as to gain the amount of energy, stored in the accelerating structure.

We will have proton-pion colliding beams when replacing the positive pions by a proton bunch of the same momentum and length. With N_p^1 protons, we have in this case

$$L_\Sigma^{\pi p} = L_\Sigma^{\pi\pi} \frac{N_p^1}{N_\pi} ,$$

and with $N_p^1 = 10^{12}$ and other parameters as above, we expect

$$L_\Sigma^{\pi p} = 3 \times 10^{28}\ \mathrm{cm}^{-2}\ \mathrm{s}^{-1} .$$

If instead of pions we inject muon bunches of the same number N_μ into the same storage ring, cooled and accelerated as already described, we have the average luminosity

$$L_\Sigma^{\mu\mu} = \frac{\zeta \dot{N}_p N_\mu}{\ell_{cool} \cdot \beta_0} \frac{p}{2m_e c} \frac{eH\tau_\mu}{2\pi m_\mu c} ,$$

where

ℓ_{cool} = the length of the cooling target equal to the β-function at this point,

β_0 = the value of the β-function in the interaction region of the storage ring.

It must be mentioned that we have made the optimistic assumption that during the collisions the muon beams have the same normalized emittance as after the ionization cooling. Assuming ℓ_{cool} = 1 cm, β_0 = 5 cm, and that the other parameters are the same as those for the pion beams, we shall have under these conditions

$$L_\Sigma^{\mu\mu} = 3 \times 10^{31}\ \mathrm{cm}^{-2}\ \mathrm{s}^{-1} .$$

5.9 Application to e^+e^- linear collider

In addition to the possibilities discussed above it will be possible to use the intense beams of large proton accelerators for exciting accelerating structures for electron-positron linear colliding beams[7], and special efforts must be made to provide a fast damping of asymmetrical modes that can be excited in the accelerating structure by proton beams.

In this case the luminosity that can be achieved is difficult to evaluate, because its value depends on how small the effective emittances of $e^\pm$ beams can be made in the collision point. If we follow the optimistic approach of Ref. 7, there will be the luminosity

$$L_\Sigma^{e^+e^-} = 10^{31}\,\mathrm{cm}^{-2}\ \mathrm{s}^{-1}$$

with $\dot{N} = 10^{13}$ pps. Also, it is very probable that the productivity of proton accelerators will still be growing.

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